

Collective stimulated Brillouin scatter

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We develop a statistical theory of stimulated Brillouin backscatter (BSBS) of a spatially and temporally partially incoherent laser beam for laser fusion relevant plasma. We find a new collective regime of BSBS which has a much larger threshold than the classical threshold of a coherent beam in long-scale-length laser fusion plasma. We identify two contributions to BSBS convective instability increment. The first is collective with intensity threshold independent of the laser correlation time and controlled by diffraction. The second is independent of diffraction, it grows with increase of the correlation time and does not have an intensity threshold. The instability threshold is inside the typical parameter region of National Ignition Facility (NIF). We also find that the bandwidth of KrF-laser-based fusion systems would be large enough to allow additional suppression of BSBS.

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Inertial confinement fusion (ICF) experiments require propagation of intense laser light through underdense plasma subject to laser-plasma instabilities which can be deleterious for achievement of thermonuclear target ignition because they can cause the loss of target symmetry, energy and hot electron production [1]. Among laser-plasma instabilities, backward stimulated Brillouin backscatter (BSBS) has long been considered a serious danger because the damping threshold of BSBS of coherent laser beams is typically several order of magnitude less than the required laser intensity $\sim 10^{15} \text{W/cm}^2$ for ICF. BSBS may result in laser energy retracing its path to the laser optical system, possibly damaging laser components [1, 2].

Theory of laser-plasma interaction (LPI) instabilities is well developed for coherent laser beam [3]. However, ICF laser beams are not coherent because temporal and spatial beam smoothing techniques are currently used to produce laser beams with short enough correlation time, T_c , and lengths to suppress speckle self-focusing. The laser intensity forms a speckle field - a random in space distribution of intensity with transverse correlation length $l_c \simeq F\lambda_0$ and longitudinal correlation length (speckle length) $L_{\text{speckle}} \simeq 7F^2\lambda_0$, where F is the optic $f/\#$ and $\lambda_0 = 2\pi/k_0$ is the wavelength (see e.g. [4, 5]). There is a long history of study of amplification in random media (see e.g [6, 7] and references there in). For small laser beam correlation time T_c , the spatial instability increment is given by a Random Phase Approximation (RPA). Beam smoothing for ICF typically has T_c much larger than the for the regime of RPA applicability. There are few examples in which the implications of laser beam spatial and temporal incoherence have been analyzed for such larger T_c . One exception is forward stimulated Brillouin scattering (FSBS). Although FSBS for a strictly coherent laser beam is a classic linear theory, we have obtained [8, 9] its dispersion relation for laser beam correlation time T_c too large for RPA relevance,

but T_c small enough to suppress single laser speckle instabilities [10]. We verified our theory of this "collective" FSBS regime with 3D simulations. Similar simulation results had been previously observed [11]. This naturally leads one to consider the possibility of a collective regime for BSBS backscatter (CBSBS). We will present 2D and 3D simulation results as evidence for such a regime, and find agreement with a simple theory that above CBSBS threshold, the spatial increment for backscatter amplitude κ_i , is well approximated by the sum of two contributions. The first is RPA-like $\propto T_c$ without intensity threshold (we neglect light wave damping). The second has a threshold as a function of laser intensity. For National Ignition Facility NIF parameters the threshold is comparable with NIF intensities. That second contribution is collective-like because it neglects speckle contributions and is only weakly dependent on T_c . CBSBS threshold is applicable for strong and weak acoustic damping coefficient ν_{ia} . The theory also provides a good quantitative prediction of the instability increment for small $\nu_{ia} \sim 0.01$ which is relevant for gold plasma near the wall of hohlraum in NIF experiments[1].

Assume that laser beam propagates in plasma with frequency ω_0 along z . The electric field \mathcal{E} is given by

$$\mathcal{E} = (1/2)e^{-i\omega_0 t} \left[E e^{ik_0 z} + B e^{-ik_0 z - i\Delta\omega t} \right] + c.c., \quad (1)$$

where $E(\mathbf{r}, z, t)$ is the envelope of laser beam and $B(\mathbf{r}, z, t)$ is the envelope of backscattered wave, $\mathbf{r} = (x, y)$, and c.c. means complex conjugated terms. Frequency shift $\Delta\omega = -2k_0 c_s$ is determined by coupling of E and B through ion-acoustic wave with phase speed c_s and wavevector $2k_0$ with plasma density fluctuation δn_e given by $\frac{\delta n_e}{n_e} = \frac{1}{2}\sigma e^{2ik_0 z + i\Delta\omega t} + c.c.$, where $\sigma(\mathbf{r}, z, t)$ is the slow envelope and n_e is the average electron density, assumed small compared to critical density, n_c . The cou-

pling of E and B to plasma density fluctuations gives

$$R_{EE}^{-1}E \equiv \left[i(c^{-1}\partial_t + \partial_z) + \frac{1}{2k_0}\nabla^2 \right] E = \frac{k_0}{4} \frac{n_e}{n_c} \sigma B, \quad (2)$$

$$R_{BB}^{-1}B \equiv \left[i(c^{-1}\partial_t - \partial_z) + \frac{1}{2k_0}\nabla^2 \right] B = \frac{k_0}{4} \frac{n_e}{n_c} \sigma^* E, \quad (3)$$

$\nabla = (\partial_x, \partial_y)$, and σ is described by the acoustic wave equation coupled to the ponderomotive force $\propto \mathcal{E}^2$ which results in the envelope equation

$$R_{\sigma\sigma}^{-1}\sigma^* \equiv [i(c_s^{-1}\partial_t + 2\nu_{ia}k_0 + \partial_z) - (4k_0)^{-1}\nabla^2]\sigma^* = -2k_0 E^* B. \quad (4)$$

The response of the slowly varying part of δn_e to the slowly varying part of the ponderomotive force, proportional to $|E|^2 + |B|^2$, responsible for self-focusing, is neglected. $\nu_{ia} = \nu_L/2k_0c_s$ is the scaled acoustic Landau damping coefficient. E and B are in thermal units (see e.g. [8]).

Assume that laser beam was made partially incoherent through induced spacial incoherence beam smoothing [15] which defines stochastic boundary conditions at $z = 0$ for the spacial Fourier transform (over \mathbf{r}) components $\hat{E}(\mathbf{k})$, of laser beam amplitude [8]:

$$\begin{aligned} \hat{E}(\mathbf{k}, z=0, t) &= |E_{\mathbf{k}}| \exp[i\phi_{\mathbf{k}}(t)], \\ \langle \exp i[\phi_{\mathbf{k}}(t) - \phi_{\mathbf{k}'}(t')] \rangle &= \delta_{\mathbf{k}\mathbf{k}'} \exp(-|t - t'|/T_c), \\ |E_{\mathbf{k}}| &= \text{const}, \quad k < k_m; \quad E_{\mathbf{k}} = 0, \quad k > k_m, \end{aligned} \quad (5)$$

chosen as the idealized "top hat" model of NIF optics [16]. Here $k_m \simeq k_0/(2F)$ and the average intensity, $\langle |E|^2 \rangle = \langle |E|^2 \rangle = I$ determines the constant.

In linear approximation, assuming $|B| \ll |E|$ so that only the laser beam is BSBS unstable, we can neglect right hand side (r.h.s.) of Eq. (2). The resulting linear equation with top hat boundary condition (5) has the exact solution as decomposition of E into Fourier series, $E(\mathbf{r}, z, t) = \sum_j E_{\mathbf{k}_j}$ with $E_{\mathbf{k}_j} \propto \exp[i(\phi_{\mathbf{k}_j}(t - z/c) + \mathbf{k}_j \cdot \mathbf{r} - \mathbf{k}_j^2 z/2k_0)]$.

Figures 1 show the increment κ_i of the spatial growth of backscattered light intensity $\langle |B|^2 \rangle \propto e^{-2\kappa_i z}$ as a function of the rescaled correlation time $\tilde{T}_c \equiv T_c k_0 c_s / 4F^2$ (note that definition is different by a factor $1/2F$ from the definition used for FSBS [8, 9]) obtained from the numerical solution of the linearized equations (2)-(4) using operator splitting method along the characteristics of E and B . Here and below we use dimensionless units with k_0/k_m^2 as the unit in z direction, $k_0/k_m^2 c_s$ is the time unit and $\mu \equiv 2\nu_{ia}k_0^2/k_m^2$. Also $\langle \dots \rangle$ means averaging over the statistics of laser beam fluctuations (5) and \tilde{I} is the scaled dimensionless laser intensity defined as $\tilde{I} = \frac{4F^2}{\nu_{ia}} \frac{n_e}{n_c} I$. Figure 1a corresponds to the 3 + 1D simulations (three spatial coordinates and t) with the boundary and initial conditions (5) in the limit $c \rightarrow \infty$ (i.e., setting c^{-1} terms in (2)-(4) to be zero. Figure 1b

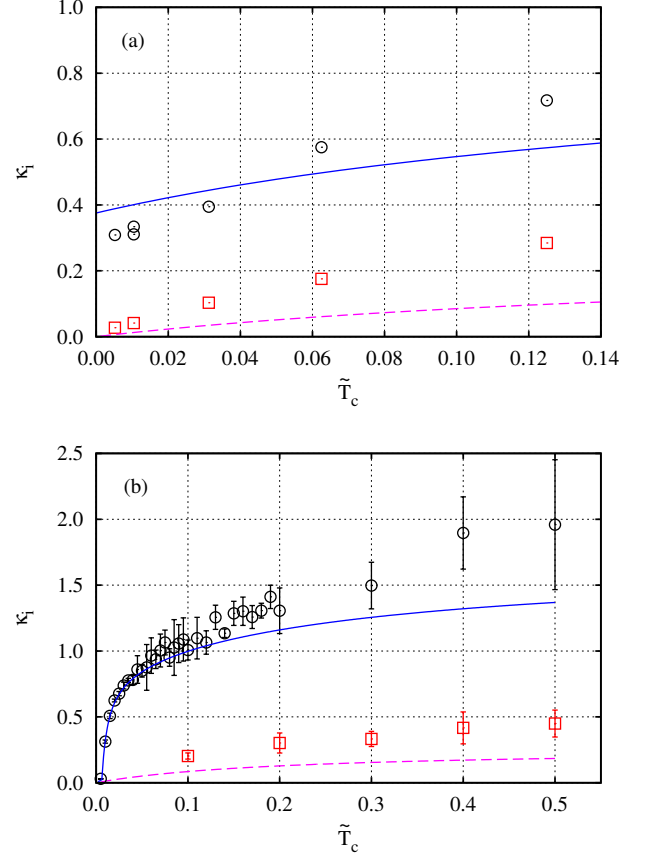


FIG. 1: Spatial increment κ_i of CBSBS obtained from numerical simulations compared with the sum of increments $\kappa_B + \kappa_\sigma$ (obtained by solving (8) and (10)). Parameters of simulations are $\nu_{ia} = 0.01$, $F = 8$. (a) 3 + 1D simulations with $c_s/c = 0$, $\tilde{I} = 2$ (circles) and $\tilde{I} = 1$ (squares). Solid and dashed lines show $\kappa_B + \kappa_\sigma$ for $\tilde{I} = 2$ and $\tilde{I} = 1$, respectively. If $\kappa_\sigma < 0$ then $\kappa_B + \kappa_\sigma$ is replaced by κ_B . (b) 2 + 1D simulations with the modified boundary conditions, $c_s/c = 1/500$, $\tilde{I} = 3$ (circles) and $\tilde{I} = 1$ (squares). Error bars are also shown. Solid and dashed lines show $\kappa_B + \kappa_\sigma$ for $\tilde{I} = 3$ and $\tilde{I} = 1$, respectively.

shows the result of 2 + 1D simulations (only 1 transverse spatial variable is taken into account) for the modified boundary condition compare with the last line in (5) as $|E_{\mathbf{k}}| = k^{1/2} \text{const}$, $k < k_m$; $E_{\mathbf{k}} = 0$, $k > k_m$ which is chosen to mimic the extra factor k in the integral over transverse direction of the full 3 + 1D problem. In that case $c/c_s \simeq 500$. E.g. for $\tilde{T}_c = 0.1$ we typically use 256 transverse Fourier modes and a discrete steps $\Delta z = 0.15$ in dimensionless units with the total length of the system $L_z = 100$ and a time step $\Delta t = \Delta z/c$. For each simulation we typically have to wait $\sim 10^6$ time steps to achieve a statistical steady state and then average over next $\sim 10^6$ time steps to find κ_i .

We now relate κ_i to the instability increments for $\langle B \rangle$ and $\langle \sigma^* \rangle$ (we designate them κ_B and κ_σ , respectively).

In general, growth rates of mean amplitudes only give a lower bound to κ_i . First we look for κ_σ . Eq. (3) is linear in B and E which implies that B can be decomposed into $B = \sum_j B_{\mathbf{k}_j}$. We approximate r.h.s. of (4) as $E^*B \simeq \sum_j E_{\mathbf{k}_j}^* B_{\mathbf{k}_j}$ so that

$$R_{\sigma\sigma}^{-1}\sigma^* = -2k_0 \sum_j E_{\mathbf{k}_j}^* B_{\mathbf{k}_j}, \quad (6)$$

which means that we neglect off-diagonal terms $E_{\mathbf{k}_j}^* B_{\mathbf{k}_{j'}}$, $j \neq j'$. Since speckles of laser field arise from interference of different Fourier modes, $j \neq j'$, we associate the off-diagonal terms with speckle contribution to BSBS [4, 12, 17]. The neglect of off-diagonal terms requires that during time T_c light travels much further than a speckle length, $L_{\text{speckle}} \ll cT_c$ and that $T_c \ll t_{\text{sat}}$, where t_{sat} is the characteristic time scale at which BSBS convective gain saturates at each speckle [13].

Eqs. (3) and (6) result in the closed expression $R_{\sigma\sigma}^{-1}\langle\sigma^*\rangle = -(k_0^2/2)(n_e/n_c)\langle E^* R_{BB} \sigma^* E \rangle$ which has the same form as the Bourret approximation [7]. We look for the solution of that expression in exponential form $B_j, \sigma^* \propto e^{i(\kappa z + \mathbf{k} \cdot \mathbf{r} - \omega t)}$, then the exponential time dependence in (5) allows to carry integrations in that expression explicitly to arrive at the following relation in dimensionless units

$$\begin{aligned} & -i\omega + \mu + i\kappa - (i/4)k^2 \\ & = 8iF^4 \frac{n_e}{n_c} \sum_{j=1}^N \frac{|E_{\mathbf{k}_j}|^2}{\omega \frac{c_s}{c} + \kappa - k_j^2 - \frac{k^2}{2} - \mathbf{k}_j \cdot \mathbf{k} + 2i \frac{c_s}{c} \frac{1}{T_c}}, \end{aligned} \quad (7)$$

where $1/k_m$ is the transverse unit of length and vectors \mathbf{k}_j span the entire top hat (5), i.e. $I = \sum_j |E_{\mathbf{k}_j}|^2$.

In the continuous limit $N \rightarrow \infty$, sum in (7) is replaced by integral which gives for the most unstable mode $\mathbf{k} = 0$:

$$-i\omega + \mu + i\kappa + i\frac{\mu}{4}\tilde{I} \ln \frac{1 - \kappa - \omega \frac{c_s}{c} - 2i \frac{c_s}{c} \frac{1}{T_c}}{-\kappa - \omega \frac{c_s}{c} - 2i \frac{c_s}{c} \frac{1}{T_c}} = 0. \quad (8)$$

The relation (8) supports the convective instability with the increment $\kappa_\sigma \equiv \text{Im}(\kappa) > 0$ only for $\tilde{I} > \tilde{I}_{\text{convthresh}}$, where $\tilde{I}_{\text{convthresh}}$ is the convective CBSBS threshold given by

$$\tilde{I}_{\text{convthresh}} = \frac{4F^2 n_e}{\nu_{ia} n_c} I_{\text{convthresh}} = 4/\pi. \quad (9)$$

In the limit $c/c_s \rightarrow \infty$, the increment κ_σ is independent of \tilde{T}_c which suggests that we refer to it as the collective instability branch. For finite but small $c_s/c \ll 1$ and $\tilde{I} > \tilde{I}_{\text{convthresh}}$ there is sharp transition of κ_σ as a function of \tilde{T}_c from 0 for $\tilde{T}_c = 0$ to \tilde{T}_c -independent value of κ_σ . That value can be obtained analytically from (8) for \tilde{I} just above the threshold as follows: $\kappa_i = \mu(\pi/4)(\tilde{I} - I_{\text{convthresh}})/(\mu\tilde{I} - 1)$.

The increment κ_B is obtained in a similar way by statistical averaging of equation (3) for $\langle B \rangle$ with σ^* from equation (4) which gives

$$-i\omega \frac{c_s}{c} + i\kappa + i\frac{\mu}{4}\tilde{I} \frac{1}{\kappa - \omega - i\mu - i\frac{1}{T_c}} = 0. \quad (10)$$

Here we neglected the contribution to $\kappa_B \equiv \text{Im}(\kappa)$ from diffraction which gives negligible correction. Equation (10) does not have a convective threshold (provided we neglect here light wave damping) while κ_B has near-linear dependence on \tilde{T}_c : $\kappa_B \simeq \mu\tilde{I}\tilde{T}_c/4$ for $\tilde{T}_c < 1/\mu$ which is typical for RPA results. It suggests that we refer κ_B as the RPA-like branch of instability.

We choose $\omega = 0.5$ in (8) and $\omega = 0$ in (10) to maximize κ_σ and κ_B , respectively. Equation (8) also predicts absolute instability for $\tilde{I} > \mu + 3\mu^{-1} + O(\mu^{-3}) + O(\tilde{T}_c^{-1}c_s/c)$, which is slightly above the coherent absolute threshold $\tilde{I} = \mu$ but here we emphasize the convective regime. Figures 1a and b show that the analytical expression $\kappa_B + \kappa_\sigma$ is a reasonable good approximation for numerical value of κ_i above the convective threshold (9) for $\tilde{T}_c \lesssim 0.1$ which is the main result of this Letter. Below this threshold the analytical and numerical results are in qualitative agreement at best but in that case we replace $\kappa_B + \kappa_\sigma$ by κ_B because $\kappa_\sigma < 0$ in that case.

The qualitative explanation why $\kappa_B + \kappa_\sigma$ is a surprisingly good approximation to κ_i is based on the following argument. First imagine that B propagates linearly and not coupled to the fluctuations of σ^* , so its source is $\sigma^*E \rightarrow \langle\sigma^*\rangle E$ in r.h.s of (3). If $\langle\sigma^*\rangle \propto e^{\kappa_\sigma z}$ grows slowly with z (i.e. if $\langle\sigma^*\rangle$ changes a little over the speckle length L_{speckle} and time T_c), then so will $\langle|B|^2\rangle$ at the rate $2\kappa_\sigma$. But if the total linear response R_{BB}^{tot} (R_{BB}^{tot} is the renormalization of bare response R_{BB} due to the coupling in r.h.s of (3)) is unstable then its growth rate gets added to κ_σ in the determination of $\langle|B|^2\rangle$ since in all theories which allow factorization of 4-point function into product of 2-point functions, $\langle B(1)B^*(2) \rangle = R_{BB}^{\text{tot}}(1,1')S(1',2')R_{BB}^{\text{tot}*}(2',2)$. Here $S(1,2) \equiv \langle\sigma^*(1)\sigma(2)\rangle\langle E(1)E^*(2) \rangle \simeq \langle\sigma^*(1)\rangle\langle\sigma(2)\rangle\langle E(1)E^*(2) \rangle$ and "1", "2" etc. mean of all spatial and temporal arguments.

The applicability conditions of the Bourret approximation used in derivation of (8) and (10) in the dimensionless units are

$$\Delta\omega_B \Delta\omega_\sigma \gg \gamma_0^2. \quad (11)$$

and $\Delta\omega_B \gg (c/c_s)|\kappa_B|$ as well as $\Delta\omega_\sigma \gg \mu$. Here γ_0 is the temporal growth rate of the spatially homogeneous solution which is given by $\gamma_0^2 = (1/4)(c/c_s)\mu\tilde{I}$. Also $\Delta\omega_\sigma = 1/\tilde{T}_c$ is the bandwidth for σ and $\Delta\omega_B$ is the effective bandwidth for B . $\Delta\omega_B$ is dominated by the diffraction in (3) which gives in the dimensionless units $\Delta\omega_B = c/c_s$. Then (11) reduces to $\tilde{T}_c \ll 4/(\mu\tilde{I})$ and $|\kappa_B| \ll 1$. Together with the condition $T_c \gg L_{\text{speckle}}/c$ used in

the derivation of (8) and assuming that $\tilde{I} \simeq \tilde{I}_{convthresh}$, it gives a double inequality $(7\pi/2)(c_s/c) \ll \tilde{T}_c \ll \pi/\mu$ which can be well satisfied for $\mu \simeq 5$, i.e. for $\nu_{ia} \simeq 0.01$ as in gold NIF plasma but not for $\mu \simeq 50$ as in low ionization number Z NIF plasma. Also $|\kappa_B| < 1$ implies that $\tilde{I} > \tilde{I}_{convthresh}$ because otherwise, below that threshold, $\kappa_B \sim -\mu$ which would contradict $|\kappa_B| < 1$. All these conditions are satisfied for $\tilde{T}_c \lesssim 1/4$ for the parameters of Figure 1 with $\tilde{I} = 2$ or $\tilde{I} = 3$ (solid lines in Figure 1) but not for $\tilde{I} = 1$ (dashed lines in Figure 1). Additionally, an estimate for $T_c \ll t_{sat}$ from the linear part of the theory of Ref. [13] results in the condition $\tilde{T}_c \ll 8\tilde{I}/\mu$ which is much less restrictive than the previous condition. These estimates are consistent with the observed agreement between $\kappa_i = \kappa_\sigma + \kappa_B$ and κ_i from simulations (filled circles in Figure 1) for \tilde{I} above the threshold (9). We conclude from Figure 1 that the applicability condition for the Bourret approximation is close to the domain of \tilde{T}_c values for which $\kappa_i = \kappa_\sigma + \kappa_B$.

For typical NIF parameters [1, 9], $F = 8$, $n_e/n_c = 0.1$, $\lambda_0 = 351\text{nm}$ and $c_s = 6 \times 10^7 \text{ cm s}^{-1}$ and the electron plasma temperature $T_e \simeq 5\text{keV}$, we obtain from (9) that $I_{convthresh} \simeq 2.2 \times 10^{14} \text{ W/cm}^2$ for gold plasma with $\nu_{ia} \simeq 0.01$, in the range of NIF single polarization intensities. So we conclude that for gold NIF plasma $I \sim I_{convthresh}$ while for low Z plasma with $\nu_{ia} \sim 0.1$ I is well below $I_{convthresh}$. Fig. 2 shows κ_i in the limit $c_s/c = 0$, $\tilde{T}_c \rightarrow 0$ from simulations, analytical result κ_σ ($\kappa_B = 0$ in that limit) and the instability increment of the coherent laser beam $\kappa_{coherent} = \mu/2 - (\mu^2 - \mu\tilde{I})^{1/2}/2$ (see e.g. [3]). It is seen that the coherent increment significantly overestimates numerical increment especially around $I_{convthresh}$. If we include the effect of finite $c_s/c = 1/500$ and finite \tilde{T}_c as in Fig. 1b then κ_i has a significant dependence on \tilde{T}_c . Current NIF 3Å beam smoothing design has $T_c \simeq 4\text{ps}$ which implies $\tilde{T} \simeq 0.15$. In that case Fig. 1b shows that there is a significant (about 5 fold) change in κ_i between $\tilde{I} = 1$ and $\tilde{I} = 3$. Similar estimate for KrF lasers ($\lambda_0 = 248\text{nm}$, $F = 8$, $T_c = 0.7\text{ps}$) gives $\tilde{T}_c = 0.04$ [ps] which results in a significant (40%) reduction of κ_i for $\tilde{I} = 3$ compare with above NIF estimate.

For practical application the threshold of BSBS is often understood as the total gain required to amplify initial thermal fluctuations up to $|B|^2 \sim |E|^2$. With such definition of threshold our results indicate that the coherent BSBS increment significantly overestimates κ_i for practical values of \tilde{T}_c as can be seen from a comparison of Figures 1b and 2.

T_c in NIF can be further reduced by self-induced temporal incoherence with collective FSBS decreasing the correlation length with beam propagation [8, 9]. For low Z plasma threshold for the collective FSBS is close to (9) [8]. As Z increases (which can be achieved by adding high Z dopant), that threshold decreases below (9) and might result in an increase of the BSBS threshold.

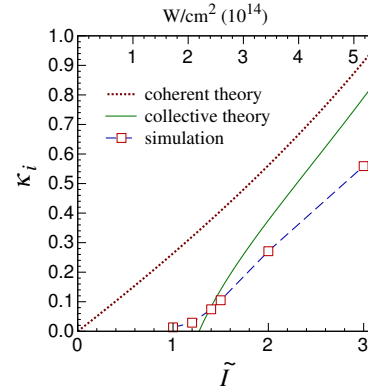


FIG. 2: κ_i vs. \tilde{I} for $\mu = 5.12$ obtained from simulations (squares connected by dashed line, $c_s/c = 0$ and limit $\tilde{T}_c \rightarrow 0$ taken by extrapolation from $\tilde{T}_c \ll 1$), analytical result κ_σ (solid curve) and coherent laser beam increment $\kappa_{coherent}$ (dotted curve). Upper grid corresponds to laser intensity in dimensional units for NIF parameters and gold plasma $T_e \simeq 5\text{keV}$, $F = 8$, $n_e/n_c = 0.1$, $\nu_{ia} = 0.01$, $\lambda_0 = 351\text{nm}$.

In conclusion, we identified a collective threshold for BSBS instability of partially incoherent laser beam for ICF relevant plasma. Above that threshold the BSBS increment κ_i is well approximated by the sum of the collective like increment κ_σ and RPA-like increment κ_B . We found that κ_i is significantly below the BSBS increment $\kappa_{coherent}$ of the coherent laser beam.

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